

Chiral Gepner Model Orientifolds

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ABSTRACT

We summarize recent progress in constructing orientifolds of Gepner models, a phenomenologically interesting class of exactly solvable string compactifications with viable gauge groups and chiral matter.

1. Introduction

The advent of D-branes has not only led to a comprehensive picture of non-perturbative aspects of string theory and its various dualities, but has also broadened our view on how our low-energy world might be embedded into string theory. In particular, the world-volume theory on D-branes canonically is a gauge theory so that string compactifications to four dimensions with D-branes in the background are interesting models for string phenomenology.

It is quite intriguing that D-branes naturally not only give rise to gauge fields but in addition provide a mechanism to yield chiral fermions. Namely, two non-trivially intersecting D-branes support a chiral fermion on each intersection point [1]. Various so-called intersecting D-brane models [2,3,4,5,6] have been studied during the last couple of years, where mostly the background geometry was chosen to be either a torus or a toroidal orbifold (see [7,8] for reviews). The most promising class of intersecting D-brane models are intersecting D6-branes on orientifolds of Type IIA string theory. In this case both the orientifold planes and the D-branes wrap various non-trivially intersecting 3-cycles of the underlying geometry and fill out the four-dimensional Minkowskian space-time.

During the last year, techniques have been developed to move beyond this limited class of toroidal orbifolds and to treat also intersecting D-brane models on more general Calabi-Yau manifolds. Namely, such models were studied at very particular symmetric points in the Calabi-Yau moduli space, where the non-linear sigma models are exactly solvable in terms of superconformal field theories known as Gepner models [10]. The internal string sector is represented by a tensor product of $N = (2, 2)$ worldsheet supersymmetric minimal models. The partition function of each of these 168 theories can be modified by simple current extensions, which yields a huge class of $N = 2$ spacetime SUSY string compactifications. The SUSY is broken down to $N = 1$ by introducing D-branes and orientifold planes, which are described in terms of boundary and crosscap states respectively. One advantage of this approach is that the CFT is exact to all orders in α' .

This is particularly important since geometrically, these models are at radii of the order of the string scale, but without relying on the intuition from classical geometry. It would be interesting to study such models also in the large radius regime [9], where, combined with non-trivial background fluxes, viable standard like models with frozen closed string moduli might be possible to achieve.

Historically, the first fully fledged Gepner model orientifolds were constructed in a case by case study during the nineties both in six [11] and four [12] uncompactified dimensions. However, during the last year these models saw a revival of interest and were studied systematically [13,14,15,16,17,18,19] (for earlier work on crosscap states in Gepner models please consult the reference list in these papers).

In this article we briefly summarize the main ingredients for the construction of Gepner model orientifolds. Due to lack of space we can explain neither any technical details nor the notation used in the formulae. Therefore, we have to refer the interested reader to the existing original literature.

2. The 1-loop amplitudes

The massless spectrum can be determined from the various one-loop open string amplitudes, which generically contain divergent tadpole contributions. The main result of the papers [14,16,18] is the computation of the general form of the Klein-bottle and Möbius amplitude for A-type simple current extended Gepner model orientifolds, where the world-sheet parity transformation Ω may be dressed by additional phases, Δ_j , and quantum symmetries, ω, ω_α , [15]. Note that via the Greene/Plesser mirror construction successive orbifolding of the A-type models yields the B-type models, which have been studied in [14].

In tree channel the Klein-bottle amplitude was shown to be expressible as the overlap

$$\widetilde{K}^A = \int_0^\infty dl \langle C | e^{-2\pi l H_{cl}} | C \rangle_A$$

with the crosscap state given in the compact form

$$\begin{aligned} |C; \Delta_j, \omega, \omega_\alpha\rangle_{NS} &= \frac{1}{\kappa_c^A} \sum_{\lambda', \mu'}^{ev} \sum_{\nu_0=0}^{\frac{K}{2}-1} \sum_{\nu_j=0}^1 \sum_{\epsilon_j=0}^1 (-1)^{\nu_0} \left(\prod_{k < l} (-1)^{\nu_k \nu_l} \right) (-1)^{\sum_j \nu_j} (-1)^{\omega \frac{s'_0}{2}} \\ e^{i\pi \sum_j \frac{\Delta_j}{k_j+2} (m'_j + (1-\epsilon_j)(k_j+2))} &\left(\prod_{\alpha} \delta_{Q_{\lambda', \mu' + (1-\epsilon_j)(\vec{k}+2), \omega_\alpha}}^{(2)} \right) \delta_{s'_0+2\nu_0+2 \sum \nu_j+2, 2\omega}^{(4)} \delta_{\sum_j \frac{1}{k_j+2} (m'_j + (1-\epsilon_j)(k_j+2)), \omega}^{(2)} \\ \prod_{j=1}^r &\left(\sigma(l'_j, m'_j, s'_j) \frac{P_{l'_j, \epsilon_j k_j}}{\sqrt{S_{l'_j, 0}}} (-1)^{\epsilon_j \frac{m'_j + s'_j}{2}} \delta_{m'_j + (1-\epsilon_j)(k_j+2), 0}^{(2)} \delta_{s'_j + 2\nu_0 + 2\nu_j + 2(1-\epsilon_j), 0}^{(4)} \right) |\lambda', \mu'\rangle_c. \end{aligned}$$

As expected the crosscap state is a sum over the crosscap Ishibashi states weighted with essentially the modular $P = T^{\frac{1}{2}} S T^2 S T^{\frac{1}{2}}$ matrix of the tensor product of the $N = 2$ minimal models subject to extra projections and non-trivial sign factors, which are derived from consistency of the modular transformation of the Möbius amplitude.

In order to cancel the tadpoles from the KB amplitude, we introduce simple current invariant Cardy-like A-type boundary states a la Recknagel/Schomerus [20]

$$\begin{aligned} |a\rangle_A &= |S_0; (L_j, M_j, S_j)_{j=1}^r\rangle_A = \frac{1}{\kappa_a^A} \sum_{\lambda', \mu'}^\beta \prod_\alpha \delta^{(1)}(Q_{\lambda', \mu'}^{(\alpha)}) \\ &(-1)^{\frac{s_0'^2}{2}} e^{-i\pi \frac{s_0' S_0}{2}} \prod_{j=1}^r \left(\frac{S_{l_j', L_j}}{\sqrt{S_{l_j', 0}}} e^{i\pi \frac{m_j' M_j}{k_j+2}} e^{-i\pi \frac{s_j' S_j}{2}} \right) |\lambda', \mu'\rangle. \end{aligned}$$

(For simple currents with fixed points these boundary states can be refined into resolved or fractional boundary states.)

The action of $\Omega_{\Delta_j, \omega, \omega_\alpha}$ on these boundary states reads

$$|S_0; (L_j, M_j, S_j)_{j=1}^r\rangle_A \rightarrow |-S_0 + 2\omega; (L_j, -M_j + 2\Delta_j, -S_j)_{j=1}^r\rangle_A.$$

Invariant boundary states carry $SO(N)$ or $SP(2N)$ gauge groups, whereas non-invariant ones have to be introduced in pairs and thus give rise to $U(N)$ gauge groups.

A boundary state is supersymmetric relative to the crosscap state if

$$\frac{S_0 - \omega}{2} - \sum_j \frac{M_j - \Delta_j}{k_j + 2} + \sum_j \frac{S_j}{2} = 0 \mod 2.$$

From the crosscap and the boundary states one can compute the Möbius strip amplitude, which together with the annulus amplitude allows one to read off the open string spectrum.

The tadpole cancellation condition can be easily determined from the contributions of the massless fields (λ, μ) to the amplitudes and takes the form

$$\text{Tad}_D(\lambda, \mu) - 4 \text{Tad}_O(\lambda, \mu) = 0.$$

Here, the contribution from the orientifold plane reads

$$\begin{aligned} \text{Tad}_O(\lambda, \mu) &= (-1)^{(1+\frac{s_0}{2})(1+\omega)} \sum_{\epsilon \in r=0}^1 e^{i\pi \sum_j \frac{\Delta_j}{k_j+2} (1-\epsilon_j)(k_j+2)} \left(\prod_{k < l} (-1)^{\epsilon_k \epsilon_l} \right) \delta_{\sum \epsilon_j, \omega + \frac{s_0}{2}}^{(2)} \\ &\left(\prod_\alpha \delta_{Q_{\lambda, \mu + (1-\tilde{\epsilon})(\vec{k}+2), \omega_\alpha}}^{(2)} \right) \prod_j \left(\sin \left[\frac{1}{2} (l_j, \epsilon_j k_j) \right] \delta_{l_j + (1-\epsilon_j)k_j, 0}^{(2)} \delta_{m_j + (1-\epsilon_j)(k_j+2), 0}^{(2)} (-1)^{\epsilon_j \frac{m_j}{2}} \right). \end{aligned}$$

Collecting all terms from the boundary states and their $\Omega_{\Delta_j, \omega, \omega_\alpha}$ images, their massless tadpoles are given by

$$\text{Tad}_D(\lambda, \mu) = \left(\prod_\alpha \delta_{Q_{\lambda, \mu}}^{(1)} \right) \sum_{a=1}^N 2 N_a \cos \left[\pi \sum_j \frac{m_j (M_j^a - \Delta_j)}{k_j+2} \right] \prod_j \sin(l_j, L_j^a).$$

As was shown in [15,16], these equations allow for solutions with unitary gauge symmetries and chiral fermions. This is in contrast to the pure B-type orientifolds, which, though easier to solve, always give rise to non-chiral matter.

3. Outlook

We have reviewed the construction of A-type Gepner model orientifolds. As a next step, it would be interesting to systematically scan all the possible Gepner models for

phenomenologically interesting examples. First encouraging results of such a search were reported in [17]. Moreover for phenomenological applications more refined data is needed and one has to determine the general form of the various terms in the $N = 1$ supersymmetric effective action. It would also be interesting to determine what happens if one moves away from the Gepner point and, related to this question, to study such orientifolds in the supergravity regime, where the recently developed methods of flux compactifications become applicable.

4. References

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